

Exponential Orthogonality Catastrophe at the Anderson Metal-Insulator Transition

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We consider the orthogonality catastrophe at the Anderson Metal-Insulator transition (AMIT). The typical overlap F between the ground state of a Fermi liquid and the one of the same system with an added potential impurity is found to decay at the AMIT exponentially with system size L as $F \sim \exp(-\langle I_A \rangle / 2) = \exp(-cL^\eta)$, where I_A is the so called Anderson integral, η is the power of multifractal intensity correlations and $\langle \dots \rangle$ denotes the ensemble average. Thus, strong disorder typically increases the sensitivity of a system to an additional impurity exponentially. We recover on the metallic side of the transition Anderson's result that fidelity F decays with a power law $F \sim L^{-q(E_F)}$ with system size L . This power increases as Fermi energy E_F approaches mobility edge E_M as $q(E_F) \sim (\frac{E_F - E_M}{E_M})^{-\nu\eta}$, where ν is the critical exponent of correlation length ξ_c . On the insulating side of the transition F is constant for system sizes exceeding localization length ξ . While these results are obtained from the mean value of I_A , giving the typical fidelity F , we find that I_A is widely, log normally, distributed with a width diverging at the AMIT. As a consequence, the mean value of fidelity F converges to one at the AMIT, in strong contrast to its typical value which converges to zero exponentially fast with system size L . This counterintuitive behavior is explained as a manifestation of multifractality at the AMIT.

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Anderson showed in Ref. 1 that the addition of a static potential impurity to a system of N fermions changes its ground-state such that the overlap between the original $\langle \psi |$ and the new ground state $\langle \psi' |$ has an upper bound,

$$\chi = F^2 = |\langle \psi | \psi' \rangle|^2 < \exp(-I_A), \quad (1)$$

where the Anderson integral I_A is for noninteracting electrons given in terms of the single particle eigenstates of the original system $|n\rangle$ and the new system $|m'\rangle$ by

$$I_A = \frac{1}{2} \sum_{\epsilon_n \leq 0, \epsilon_{m'} > 0} |\langle n | m' \rangle|^2. \quad (2)$$

If the added impurity is short ranged and of strength λ , Anderson found for a clean metal $I_A = 2\pi^2\lambda^2 \ln N$, diverging with the number of fermions N , so that F , also called fidelity, decays with a power law with N , leading to the so called orthogonality catastrophe (AOC). This implies that the local perturbation connects the system to a macroscopic number of excited states which has important consequences like the singularities in the X-Ray absorption and emission of metals[2]. Furthermore, the zero bias anomaly in disordered metals[3] and anomalies in the tunneling density of states in quantum Hall systems[4] are related to the AOC. The concept of fidelity can be generalised to any parametric perturbation of a system and be used to characterise quantum phase transitions [5]. The AOC has been explored in mesoscopic systems[6, 7]. With the advent of engineered many-body systems in ensembles of ultracold atoms it is possible to study nonequilibrium quantum dynamics of such systems in a controlled way so that consequences of parameter changes become measurable directly[8].

An intriguing question is, if the system becomes less or more sensitive to the addition of another impurity if it already

contains a finite density of impurities. Gefen et al. showed in Ref. 9 that in a weakly disordered metal the average value $\langle I_A \rangle$ scales with $\ln N$ when the potential of the added impurity potential is short ranged. That result is valid to leading order in $1/g$, where $g = k_F l_e \gg 1$ is a measure of disorder strength with mean free path l_e . Numerical results [9] show that in a 2-dimensional disordered system I_A increases as the disorder strength increases until it is so strong that localization length ξ is smaller than system size L . Beyond that, I_A decreases as ξ decreases with increasing disorder. Thus, the addition of an impurity changes the ground state of weakly disordered systems more strongly than the one of a clean system. Only at strong disorder when the fermions are localized, its sensitivity to a potential change decreases again.

Here, we aim to derive this behavior analytically in order to find out how fidelity F changes when tuning disorder strength or energy. In systems close to the Anderson metal-insulator transition (AMIT) we can use the fact that the single particle wavefunctions at the AMIT are multifractal [10] and power law correlated [11–13]. We also obtain analytical results for noncritical 2-dimensional disordered systems, which confirm the numerical calculations of Ref. 9. For a short range impurity of strength λ , located at position \mathbf{x} , I_A can be expressed in terms of the local intensities of the unperturbed Eigenstates $|n\rangle$ with Eigenenergies E_n , [1, 9]

$$I_A = \frac{(2\pi\lambda)^2}{2\rho^2} \sum_{E_n < E_F} \sum_{E_m > E_F} \frac{|\psi_n(\mathbf{x})|^2 |\psi_m(\mathbf{x})|^2}{(E_n - E_m)^2}, \quad (3)$$

where ρ is the mean density of states. The correlation function of the intensities associated to two energy levels distant by

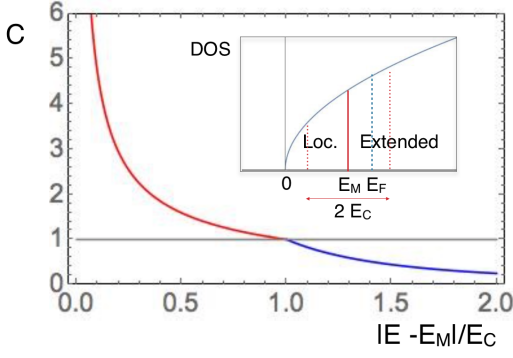


FIG. 1: Correlation function of intensities C as function of their energy difference $|E - E_M|$ in a disordered system with $\eta = 2$, $d = 2$. Inset: Density of states (DOS) as function of energy E with a transition between localized and extended states at mobility edge E_M . Intensities are correlated within an energy window $2E_C$ around E_M .

$\omega_{nm} = E_n - E_m$ is given by [12–14]

$$C(\omega_{nm} = E_n - E_m) = L^d \int d^d r \langle |\psi_n(\mathbf{r})|^2 |\psi_m(\mathbf{r})|^2 \rangle$$

$$= \begin{cases} \left(\frac{E_c}{\text{Max}(|\omega_{nm}|, \Delta)} \right)^{\eta/d}, & 0 < |\omega_{nm}| < E_c, \\ \left(\frac{E_c}{|\omega_{nm}|} \right)^2, & |\omega_{nm}| > E_c, \end{cases} \quad (4)$$

when $E_n \leq E_M$ and $E_m \geq E_M$ or vice versa [15]. $\Delta = 1/(\rho L^d)$ is the average level spacing. Here, $\eta = 2(\alpha_0 - d)$, with multifractality parameter α_0 and d the dimension. The correlation energy E_c is a macroscopic energy of order of elastic scattering rate $1/\tau$. For $|\omega_{nm}| < E_c$ correlations are enhanced in comparison to the plane-wave limit $C_{nm} = 1$, see Fig. 1, where we set one of the energies at the mobility edge $E_n = E_M$ and the other at $E_m = E$. Note, that for $|\omega_{nm}| > E_c$ it decays below 1. This anticorrelation ensures that the total intensity at a position \mathbf{x} is normalised: A dip in intensity at one energy implies an enhancement of intensity at another energy and vice versa.

Mean Value of the Anderson Integral.— Inserting Eq. (4) into Eq. (3), the average mean value of I_A is

$$\langle I_A \rangle = \frac{(2\pi\lambda)^2}{2} \iint_{\epsilon < -\Delta/2, \epsilon' > \Delta/2} d\epsilon d\epsilon' \frac{C_{\epsilon, \epsilon'}}{(\epsilon - \epsilon')^2}. \quad (5)$$

This gives the geometrical average of F , $\exp(\langle \ln F \rangle) = \exp(-\langle I_A \rangle/2)$. At the AMIT we get with Eq. (4)

$$\langle I_A \rangle|_{E_F = E_M} = \frac{2(\pi\lambda)^2}{\gamma(1+\gamma)} \left(\frac{E_c}{\Delta} \right)^\gamma, \quad (6)$$

diverging with number of particles $N = E_F/\Delta$ with a power law, $\gamma = \eta/d$.

As the Fermi energy is moved into the insulating regime $E_F < E_M$, there remain multifractal correlations, Eq. (4), but the integral is now cut off at local level spacing $\Delta_\xi =$

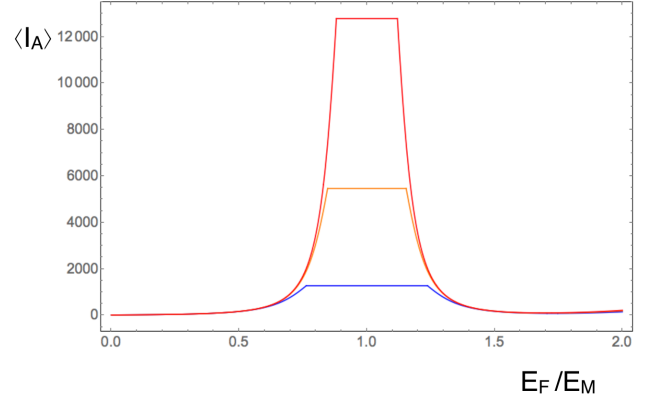


FIG. 2: The average Anderson integral $\langle I_A \rangle$ as function of Fermi energy E_F in units of mobility edge E_M for a 3-dim. disordered system, Eqs. (6,7,8), for various system sizes L (in units of microscopic length a_c), $L = 10$ (blue), $L = 20$ (orange) and $L = 50$ (red).

$1/(\rho \xi^d)$, since there is local level repulsion [15]. This yields

$$\langle I_A \rangle|_{E_F < E_M} = \frac{2(\pi\lambda)^2}{\gamma(1+\gamma)} \left(\frac{E_c}{\Delta_\xi} \right)^\gamma, \quad (7)$$

independent of the number of particles N .

In the metallic regime $E_F > E_M$ all wave functions are extended. On length scales smaller than correlation length ξ_c multifractal fluctuations still occur and there are power-law correlations in energy Eq. (4). The energy difference $|\omega_{nm}|$ is for $\epsilon_n < \epsilon_m$ substituted by $\text{Max}[|\omega_{nm}|, \Delta_{\xi_{cn}}]$, where $\Delta_{\xi_{cn}} = E_c (\xi_{cn}/a_c)^{-d}$. ξ_{cn} is the correlation length at energy E_n and a_c is a small length scale defined by $E_c = 1/(\rho a_c^d)$ [15]. For $\omega_{nm} < E_c$, correlations are enhanced in comparison to plane-wave limit $C_{nm} = 1$, yielding

$$\langle I_A \rangle|_{E_F > E_M} = 2(\pi\lambda)^2 \left(\frac{E_c}{\Delta_{\xi_c}} \right)^\gamma \left(\frac{1}{\gamma(1+\gamma)} + \ln \frac{N}{N_{\xi_c}} \right), \quad (8)$$

diverging logarithmically with number of electrons N in agreement with Anderson's result for a metal, which is recovered exactly far away from the MIT, where $\Delta_{\xi_c} \rightarrow E_c$.

In Fig. 2 we plot the first moment of the Anderson Integral $\langle I_A \rangle$ as function of Fermi energy E_F in units of mobility Edge E_M , Eqs. (6,7,8), for various system sizes L .

Anderson Integral of 2D Disordered Electron Systems.— In 2D disordered electron systems without spin-orbit interaction, without strong magnetic field and strong interactions all states are localized. Therefore, the Anderson integral is given by Eq. (7) with localization length $\xi_{2D} = g \exp(\pi g)$ (in units of the smallest microscopic length scale), where $g = \epsilon_F \tau$. In 2D weakly disordered systems, $g \gg 1$, there are weak logarithmic correlations of the intensity at different energies which can to leading order in $1/g$ be written as power law correlations with power $\eta_{2D} = 2/(\pi g)$. Substituting ξ_{2D} and $\gamma_{2D} = 1/(\pi g)$ into Eq. (7), with $\Delta_\xi = D/\xi^2$ and $\Delta = D/L^2$, where D is the band width, we get the Ander-

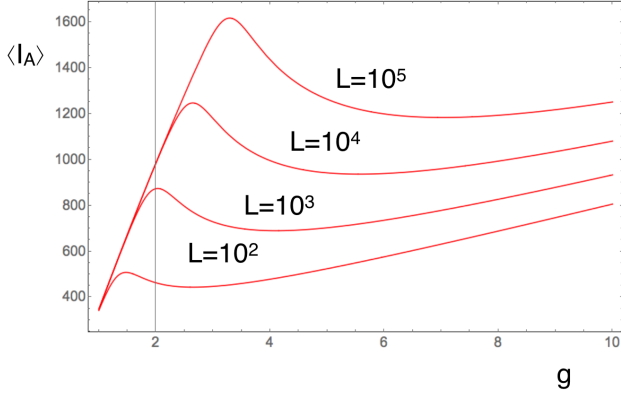


FIG. 3: The average Anderson integral $\langle I_A \rangle$ as function of disorder parameter $g = E_F \tau$ for a 2-dimensional disordered system, as obtained by substituting $\xi_{2D} = g \exp(\pi g)$ and $\gamma_{2D} = 1/(\pi g)$ into Eq. (7), with $\Delta_\xi = D/\xi^2$ and $\Delta = D/L^2$ for various system sizes L .

son integral for 2D disordered systems as function of disorder parameter g and system size L as plotted in Fig. 3. We thus confirm analytically the nonmonotonous dependence of $\langle I_A \rangle$ as function of disorder strength as observed numerically in 2D disordered systems in Ref. 9. We find that the maximal value of I_A increases with system size logarithmically. When the localization length exceeds the system size $\xi > L$ we find $\langle I_A \rangle = 2\pi(\pi\lambda)^2 g L^2/(\pi g)$. Expanding in $1/g$ we recover to leading order Anderson's result, $\langle I_A \rangle = \text{const.} + 2(\pi\lambda)^2 \ln L^2$. Note that the logarithmic divergence is independent of the disorder parameter g and coincides exactly with Anderson's result for a clean metal. In the opposite limit, when localization length ξ is smaller than systems size L , we find $\langle I_A \rangle = 2\pi(\pi\lambda)^2 g \exp(2 + 2 \ln g/(\pi g))$.

Distribution Function of the Anderson Integral.— Having obtained that the average Anderson integral diverges with the system size L at the AMIT more strongly than in a clean or weakly disordered metal, we may ask how widely I_A is distributed. Since it is a functional of the local density of states at the position where the additional impurity has been placed, Eq. (3), it is expected to be widely distributed. The distribution function of $I_A(\mathbf{x})$ as obtained by placing the impurity in different ensembles with same disorder strength is

$$P(I_A) = \int \prod_l d\alpha_l P(\{\alpha_l\}) \delta(I_A - I[\{\alpha_l\}]), \quad (9)$$

where $I[\{\alpha_l\}]$ is defined by the right side of Eq. (3). Following the strategy recently used in the derivation of the distribution function of Kondo temperatures T_K at the AMIT[15], we replace the correlated distribution function of all intensities $P(\{\alpha_l\})$ by a product of pairwise joint distribution functions of $|\psi_n(\mathbf{r}_i)|^2$ and $|\psi_m(\mathbf{r}_i)|^2$, in accordance with the correlation function C_{nm} , Eq. (4). Thereby, we can derive the conditional intensity of a state at energy E_l , given that the intensity at mo-

bility edge E_M is $|\psi_M(\mathbf{r})|^2 = L^{-\alpha}$, [15]

$$L^d \langle |\psi_l(\mathbf{r})|^2 \rangle_{|\psi_M(\mathbf{r})|^2 = L^{-\alpha}} = \left| \frac{E_l - E_M}{E_c} \right|^{r_\alpha}, \quad (10)$$

where the power is given by $r_\alpha = \frac{\alpha - \alpha_0}{d} - \frac{\eta}{2d} g_{lM}$. When E_l is located away from mobility edge E_M , the coefficient $g_{lM} = \ln |(E_l - E_M)/E_c| / (d \ln L)$ vanishes for $L \rightarrow \infty$. Close to E_M it saturates: $g_{lM}|_{E_l \rightarrow E_M} \rightarrow -1$ and Eq. (10) reduces to $L^{d-\alpha}$, the local intensity at E_M relative to the intensity of an extended state L^{-d} [15]. At positions where the local intensity at the mobility edge is small, corresponding to $\alpha > \alpha_0$, it is suppressed within an energy range of order E_c around E_M , forming *local pseudogaps* with power $r_\alpha = \frac{\alpha - \alpha_0}{d}$. When the intensity at E_M is larger than its typical value, $\alpha < \alpha_0$, there are *local power law divergencies* and the local density of states is enhanced within energy range $2E_c$ around E_M , increasing as a power law when E_l approaches E_M .

Next, we can find the Anderson integral at a position \mathbf{x} , when the intensity of the state at the mobility edge $E_n = E_M$ at that position is fixed to $|\psi_n(\mathbf{x})|^2 = L^{-\alpha}$,

$$\langle I_A \rangle_\alpha = 2(\pi\lambda)^2 \iint_{\epsilon < -\Delta/2, \epsilon' > \Delta/2} d\epsilon d\epsilon' \frac{1}{(\epsilon - \epsilon')^2} \left| \frac{\epsilon}{E_c} \right|^{r_\alpha} \left| \frac{\epsilon'}{E_c} \right|^{r_\alpha}. \quad (11)$$

which yields

$$\langle I_A \rangle_\alpha = \frac{2(\pi\lambda)^2}{\gamma(1+\gamma)} \left(\frac{E_c}{D} \right)^\gamma L^{-2(\alpha-d)}. \quad (12)$$

Inverting this equation and inserting the result into the distribution function of α , $P(\alpha) = \exp[-\ln L(\alpha - \alpha_0)^2/(2\eta)]$ we get with Eq. (9),

$$P(I_A) = \frac{1}{\sqrt{8\pi\eta \ln L}} \frac{1}{I_A} e^{-\frac{(\ln \frac{I_A}{\langle I_A \rangle} + 2\eta \ln L)^2}{8\eta \ln L}}, \quad (13)$$

where $\langle I_A \rangle$ is the first moment, Eq. (6). Thus, the Anderson integral is widely, log-normally distributed with a width which increases with system size L logarithmically.

If the Fermi energy is in the insulating regime, $E_F < E_M$, there is multifractality on length scales smaller than localization length ξ and the intensity scales with ξ , $|\psi_l(\mathbf{x})|^2 = \xi^{-\alpha_l}$. Thus, we find the distribution function of I_A by replacing the system size L by ξ in the above derivation at the MIT yielding Eq. (13), where L is replaced by ξ .

On the metallic side of the transition all wave functions are extended and their intensities scale as $|\psi|^2 \sim L^{-d}$. On length scales smaller than correlation length ξ_c multifractal fluctuations of the wave function intensity occur as long as ξ_c is larger than the microscopic length scale a_c . [12, 14] In the metallic phase, moments of intensity scale with ξ_c as $L^{dq} \langle |\psi|^{2q} \rangle \sim \xi_c^{(d-d_q)(q-1)}$ [14–16]. Therefore, we define α in the metal as $L^d |\psi_l(\mathbf{r})|^2 = \xi_{cl}^{d-\alpha_l}$, where ξ_{cl} is the correlation length of state l . As the MIT is approached ξ_c diverges and is replaced by system size L , so that α crosses over to the definition used above at the MIT. It has to a good approximation the

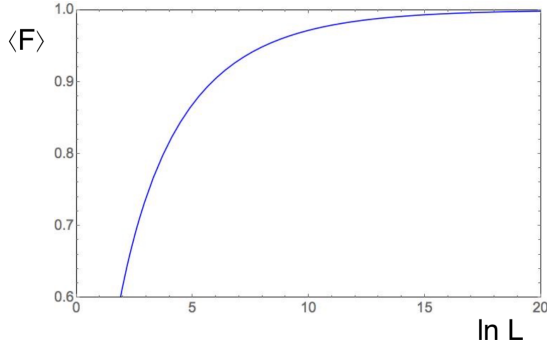


FIG. 4: The average fidelity $\langle F \rangle$ as function of the logarithm of system size L at the 3D AMIT for $\eta = 2$, $\lambda = 1$.

Gaussian distribution, $P(\alpha_l) \sim \exp[-\ln \xi_{cl}(\alpha_l - \alpha_0)^2/(2\eta)]$ as confirmed numerically in Ref. 15. Therefore, in deriving the distribution function of I_A , we can follow the strategy used at the MIT, deriving first the value of I_A when averaged over all pair correlations, given that α at the Fermi energy is fixed,

$$\langle I_A \rangle_\alpha = 2(\pi\lambda)^2 c \left(\frac{1}{\gamma(1+\gamma)} + \ln \frac{N}{N_{\xi_c}} \right) \xi_c^{-2(\alpha-d)}, \quad (14)$$

where $c = (E_c/D)^\gamma$. Inverting this equation and inserting it into the distribution function of α we get with Eq. (9) the distribution of I_A in the metallic regime, Eq. (13), replacing L by ξ_c and $\langle I_A \rangle$ by the first moment Eq. (8).

We note that the first moment of I_A yields the geometrical average of the fidelity $\exp(\langle \ln F \rangle) = \exp(-\langle I_A \rangle/2)$ giving a typical value of F . So far, we have not yet obtained the average fidelity $\langle F \rangle = \langle \exp(-I_A/2) \rangle$, since its calculation requires the knowledge of all moments of I_A [18]. Using the distribution function as obtained in the pair approximation above, Eq. (13), we find at the AMIT

$$\langle F \rangle = \frac{1}{\sqrt{8\pi\eta \ln L}} \int_{-\infty}^{\infty} dx e^{-\langle I_A \rangle L^{-\eta} \exp(x)} e^{-\frac{x^2}{8\eta \ln L}}. \quad (15)$$

An expansion in moments of I_A gives a divergent series. A saddle point approximation to the integral in Eq. (15), yields $\langle F \rangle|_{L \rightarrow \infty} = 1$, as confirmed by numerical integration, Fig. 4, where we plot $\langle F \rangle$, Eq. (15), as function of $\ln L$.

We conclude that, typically, the Anderson orthogonality catastrophe becomes exponentially enhanced at the AMIT. The typical fidelity decays exponentially with system size L as $F \sim \exp(-cL^\eta)$, where η is the power of multifractal intensity correlations. On the metallic side of the transition we recover Anderson's result that the typical overlap decays with a power law $F \sim L^{-q(E_F)}$. The power increases as Fermi energy E_F approaches mobility edge E_M like $q(E_F) \sim (\frac{E_F - E_M}{E_M})^{-\nu\eta}$, where ν is the critical exponent of correlation length ξ_c . On the insulating side of the transition the typical value of F approaches a constant for L exceeding localization length ξ . While these results for F were obtained with the mean values of Anderson integral I_A , we

derive also its distribution and find that I_A is widely, log normally, distributed with a width which diverges at the AMIT. Surprisingly, we find that the average fidelity converges at the AMIT to $F = 1$ as the system size L is sent to infinity. This is a consequence of multifractality: placing the additional short ranged impurity randomly in the sample, the fidelity is typically exponentially small. Averaging the fidelity, there is a large weight on positions where the wave function intensity is reduced and where the local density of states has a local pseudogap. At these positions the additional impurity has no effect, so that $F = 1$. As a consequence, the average fidelity is $\langle F \rangle \rightarrow 1$, while typically $\exp(\langle \ln F \rangle) \rightarrow 0$ at the AMIT in the infinite volume limit. Building on these results we can next employ the same strategy to study experimental consequences like singularities in X-Ray absorption and emission [2] and the zero bias anomaly at the MIT in doped semiconductors[17]. We also plan to extend this approach to explore the effect of more extended impurities[9], other parametric perturbations [5] and to the study of nonequilibrium quantum dynamics of disordered systems which can be studied in synthetic many-body systems with controlled disorder in ensembles of ultracold atoms.

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